

Enhanced Production of Δ and $\Sigma(1385)$ Resonances

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Yields of $\Delta(1230)$, $\Sigma(1385)$ resonances produced in heavy ion collisions are studied within the framework of a kinetic master equation. The time evolution is driven by the process $\Delta \leftrightarrow N\pi$, $\Sigma(1385) \leftrightarrow \Lambda\pi$. We obtain resonance yield both below and above chemical equilibrium, depending on initial hadronization condition and separation of kinetic and chemical freeze-out.

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Hadron resonances are produced copiously in the quark-gluon plasma (QGP) fireball break up into hadrons (hadronization, chemical freeze-out) e.g. at RHIC [1, 2, 3, 4]. Within the statistical hadronization model (SHM) approach [5, 6], the initial yields are described by chemical fugacities Υ , and hadronization temperature T . The production of heavy resonances is suppressed exponentially in m/T . Once formed, resonances decay. If this occurs inside matter, detailed balance requires also production of resonances, called ‘regeneration’ and/or ‘back-reaction’.

If the chemical freeze-out occurs much earlier than thermal, the initially produced resonances are practically invisible due to rescattering of decay products [7]. The observed yield of resonances is fixed by the physical conditions prevailing at the final breakup of the fireball, at which time last scattering occurs, this is the ‘kinetic freeze-out’. The present work addresses two questions: a) how observable resonance yield depends on the difference between chemical freeze-out temperature (e.g. point of hadronization of QGP) and the kinetic freeze-out temperature;

b) how this yield depends on the degree of initial chemical non-equilibrium at hadronization.

One can see this work as an effort to improve on the concept of chemical freeze-out for the case of resonances: given the relatively fast reactions their yield remains sensitive to the conditions prevailing between chemical and thermal freeze-out, even if this time is just 1 fm/c.

In order to describe evolution of the resonance abundance one can perform a microscopic transport simulation of the expanding system. In this approach the regeneration of resonances was previously studied by Bleicher and collaborators [8, 9, 10]. There are many detailed features of particle interactions to resolve in a microscopic model description and thus it seems appropriate to simplify the situation. We study resonance decay and regeneration using the momentum integrated population master equations, and assuming hydrodynamic expansion inspired model of fireball dynamics with conserved entropy content. In all our considerations we presume that the yield of pions π is so large that we can assume it to be unaffected by any of the reactions we consider, thus we fix pion yield in terms of fugacity and temperature

values. We do not consider all resonances which decay into resonances as does the SHARE2 SHM program [6], thus we will correct the final yields by an estimate of this effect comparing our initial resonance yield to SHARE2 results.

For the ‘fast’ baryon resonances considered here we keep the sum yields constant:

$$\begin{aligned}\Delta + N &= \Delta_0 + N_0 \equiv N_0^{\text{tot}} = \text{Const.}, \\ \Sigma(1385) + \Lambda &= \Sigma_0(1385) + \Lambda_0 \equiv \Lambda_0^{\text{tot}} = \text{Const.}\end{aligned}\quad (1)$$

The baryon annihilation, strangeness exchange such as $N + K \leftrightarrow \Lambda + \pi$ reactions, and population exchanges with higher resonances are assumed not to have a material impact within the time scale during which the temperature drops from chemical to kinetic freeze-out condition.

The experimentally observable hyperon yield appearing in our final result is

$$\Lambda_{\text{tot}} = \Sigma(1385) + \Lambda + \Sigma^0(1193) + Y^* \quad (2)$$

due to experimentally inseparable $\Sigma^0(1193) \rightarrow \gamma + \Lambda$ decay and the decay of further hyperon resonances Y^* . Similarly, when we refer to N_{tot} we include baryon resonances in the count.

In the following we will be referring explicitly to the Δ yield governed by $c\tau_{\Delta} \equiv 1/\Gamma_{\Delta} = 1.67$ fm. All equations apply equally to $\Sigma(1385)$ yield (partial decay width $\Gamma_{\Sigma \rightarrow \Lambda} \simeq 35$ MeV) and we will compare our results with experiment for this case. We note that even though the $\Sigma(1385)$ decay width is much smaller than Γ_{Δ} , the number of reaction channels and particle densities available lead to a significant effect for $\Sigma(1385)$, comparable to our finding for Δ .

The evolution in time of the Δ (or $\Sigma(1385)$) resonance yield is described by the process of resonance formation in scattering, less natural decay:

$$\frac{1}{V} \frac{dN_{\Delta}}{dt} = \frac{dW_{N\pi \rightarrow \Delta}}{dV dt} - \frac{dW_{\Delta \rightarrow N\pi}}{dV dt}, \quad (3)$$

where N_{Δ} is multiplicity of Δ resonances, $dW_{N\pi \rightarrow \Delta}/dV dt$ and $dW_{\Delta \rightarrow N\pi}/dV dt$ are invariant rates (per unit volume and time) for Δ production and decay respectively. Allowing for Fermi-blocking and

Bose enhancement in the final state, the two in-matter rates are:

$$\begin{aligned} \frac{dW_{\Delta \rightarrow N\pi}}{dVdt} &= \frac{g_\Delta}{(2\pi)^3} \int \frac{d^3 p_\Delta}{2E_\Delta} f_\Delta \int \frac{d^3 p_N}{2E_N (2\pi)^3} (1 - f_N) \\ &\times \int \frac{d^3 p_\pi}{2E_\pi (2\pi)^3} (1 + f_\pi) (2\pi)^4 \delta^4(p_N + p_\pi - p_\Delta) \\ &\times \frac{1}{g_\Delta} \sum_{spin} |\langle p_\Delta | M | p_N p_\pi \rangle|^2, \end{aligned} \quad (4)$$

$$\begin{aligned} \frac{dW_{\pi N \rightarrow \Delta}}{dVdt} &= \frac{g_N}{(2\pi)^3} \int \frac{d^3 p_N}{2E_N} f_N \frac{g_\pi}{(2\pi)^3} \int \frac{d^3 p_\pi}{2E_\pi} f_\pi \\ &\times \int \frac{d^3 p_\Delta}{2E_\Delta (2\pi)^3} (1 - f_\Delta) (2\pi)^4 \delta^4(p_N + p_\pi - p_\Delta) \\ &\times \frac{1}{g_N g_\pi} \sum_{spin} |\langle p_N p_\pi | M | p_\Delta \rangle|^2. \end{aligned} \quad (5)$$

where $g_i, i = \pi, N, \Delta$ is particles degeneracy. The distribution functions for π, N, Δ are

$$f_\pi = \frac{1}{\Upsilon_\pi^{-1} e^{u \cdot p_\pi / T} - 1}, \quad (6)$$

$$f_j = \frac{1}{\Upsilon_j^{-1} e^{u \cdot p_j / T} + 1}, \quad j = N, \Delta. \quad (7)$$

Here Υ_i is particles fugacity, and $u \cdot p_i = E_i$, for $u^\mu = (1, \vec{0})$ in the rest frame of the heat bath where $d^4 p \delta_0(p_i^2 - m_i^2) \rightarrow d^3 p_i / E_i$ for each particle. Hence, Eq.(4) and Eq.(5) are Lorentz invariant, as presented these rates can be evaluated in any convenient frame of reference. Normally, this is the frame co-moving with the thermal volume element.

Eqs.(6), (7) satisfy the usual relations for the Fermi and Bose distributions:

$$1 - f_j = \Upsilon_j^{-1} e^{u \cdot p_j} f_j, \quad j = \Delta, N \quad (8)$$

$$1 + f_\pi = \Upsilon_\pi^{-1} e^{u \cdot p_\pi} f_\pi. \quad (9)$$

Using these equations, we obtain for decay rate Eq.(4)

$$\begin{aligned} \frac{dW_{\Delta \rightarrow N\pi}}{dVdt} &= \Upsilon_N^{-1} \Upsilon_\pi^{-1} \frac{1}{(2\pi)^6} \int \frac{d^3 p_\Delta}{2E_\Delta} \int \frac{d^3 p_N}{2E_N} \\ &\times \int \frac{d^3 p_\pi}{2E_\pi (2\pi)^3} (2\pi)^4 \delta^4(p_N + p_\pi - p_\Delta) \\ &\times \sum_{spin} |\langle p_\Delta | M | p_N p_\pi \rangle|^2 f_\Delta f_N f_\pi \exp(u \cdot p_\Delta / T). \end{aligned} \quad (10)$$

where in the last exponent we replaced $p_N + p_\pi$ by p_Δ given the energy-momentum conservation 4-delta function. We can perform a similar simplification in Eq.(5). Then, observing that due to the time reversal symmetry,

$$|\langle p_1 | M | p_2 p_3 \rangle|^2 = |\langle p_2 p_3 | M | p_1 \rangle|^2$$

we find a detailed balance relation between the production and decay rate:

$$\Upsilon_\Delta \frac{dW_{N\pi \rightarrow \Delta}}{dVdt} = \Upsilon_N \Upsilon_\pi \frac{dW_{\Delta \rightarrow N\pi}}{dVdt}. \quad (11)$$

The master equation, Eq.(3), can now be cast into the form:

$$\frac{1}{V} \frac{dN_\Delta}{dt} = \left(\frac{\Upsilon_\pi \Upsilon_N}{\Upsilon_\Delta} - 1 \right) \frac{dW_{\Delta \rightarrow N\pi}}{dVdt}. \quad (12)$$

This is a rather intuitive and simple result, yet only recently the $1 \leftrightarrow 2$ population master equations have been considered [11]. Equation (12) implies for $dN_\Delta/dt = 0$ the chemical equilibrium condition:

$$\Upsilon_\pi^{\text{eq}} \Upsilon_N^{\text{eq}} = \Upsilon_\Delta^{\text{eq}}. \quad (13)$$

This equation is solved by the global chemical equilibrium $\Upsilon_\pi^{\text{eq}} = \Upsilon_N^{\text{eq}} = \Upsilon_\Delta^{\text{eq}} = 1$. However, there are also other, transient, equilibrium states possible, given a prescribed value of e.g. the background pion abundance, $\Upsilon_\pi^{\text{eq}} \neq 1$. When the initial state is formed away from transient equilibrium condition, we recognize that for $\Upsilon_\Delta < \Upsilon_\pi \Upsilon_N$ the Δ production is dominant, and conversely, for $\Upsilon_\Delta > \Upsilon_\pi \Upsilon_N$ the Δ decay dominates.

We now introduce into the population master equation (12) the effective lifespan, τ_Δ aiming to find an equation similar to classic radioactive decay population equation. We define the in medium Δ -lifespan to be:

$$\tau_\Delta \equiv \frac{\Upsilon_\Delta}{V} \frac{dN_\Delta/d\Upsilon_\Delta}{dW_{\Delta \rightarrow N\pi}/dVdt}. \quad (14)$$

We recognize that in the Boltzmann limit this corresponds to the ratio of equilibrium yield to the rate per unit time at which the equilibrium is approached. We obtain for Eq.(12):

$$\frac{dN_\Delta}{dt} = (\Upsilon_\pi \Upsilon_N - \Upsilon_\Delta) \frac{dN_\Delta}{d\Upsilon} \frac{1}{\tau_\Delta}. \quad (15)$$

In case that the ambient temperature does not vary with time, and thus only populations evolve due to change in fugacities, we have $dN/dt = dN/d\Upsilon d\Upsilon/dt$ and the following dynamical equation for the fugacity arises:

$$\tau_\Delta \frac{d\Upsilon_\Delta}{dt} = (\Upsilon_\pi \Upsilon_N - \Upsilon_\Delta). \quad (16)$$

This is ‘classical’ population equation form where the fugacity plays the role of the classical densities. When the dynamical values of $\Upsilon_i(t)$ are used in the quantum Bose/Fermi distributions, the effects of blocking, and stimulated emission are explicit.

If we instead were to introduce the lifespan by $\tilde{\tau}_\Delta \equiv (N_\Delta/V)/(dW_{\Delta \rightarrow N\pi}/dVdt)$, this implies for all particles (Bose, Fermi, Boltzmann) the classical population equation, e.g. $dN_\Delta/dt = (\Upsilon_\pi \Upsilon_N / \Upsilon_\Delta - 1) N_\Delta / \tilde{\tau}_\Delta$, and the quantum effects are now hidden in the definition of $\tilde{\tau}$.

Both definitions coincide for the case of a dilute system, and differ most for dense systems. In the limit of very dilute, vacuum system, the relaxation time is the same as the lifespan of the particles. The computed yields of particles as function of time are not dependent on the finesse of the relaxation time definition.

We now set up for semi-analytical solution of master equation (15). For multiplicities Δ and N considering the small yield and $m \gg T$ we will use the Boltzmann distribution:

$$\frac{N_\Delta}{V} = \Upsilon_\Delta \frac{T^3}{2\pi^2} g_\Delta x_\Delta^2 K_2(x_\Delta), \quad (17)$$

$$\frac{N_N}{V} = \Upsilon_N \frac{T^3}{2\pi^2} g_N x_N^2 K_2(x_N), \quad (18)$$

where $x_{\Delta,N} = m_{\Delta,N}/T$, $K_2(x)$ is Bessel function. Considering that fugacities, temperature and volume vary in time, we rewrite the left hand side of Eq.(15):

$$\frac{dN_\Delta}{N_\Delta d\tau} = \frac{d\Upsilon_\Delta}{\Upsilon_\Delta d\tau} + \frac{d \ln(x_\Delta^2 K_2(x_\Delta))}{dT} \dot{T} + \frac{d(VT^3)}{VT^3} \frac{1}{d\tau}. \quad (19)$$

We changed from t to τ to make explicit the fact that we work in fluid-element co-moving frame and thus do not consider the effect of flow on the volume time dependence.

Combining Eq.(15) with Eq.(19) we obtain

$$\frac{d\Upsilon_\Delta}{d\tau} = (\Upsilon_\pi \Upsilon_N - \Upsilon_\Delta) \frac{1}{\tau_\Delta} + \Upsilon_\Delta \frac{1}{\tau_T} + \Upsilon_\Delta \frac{1}{\tau_S}, \quad (20)$$

$$\frac{1}{\tau_T} = -\frac{d \ln(x_\Delta^2 K_2(x_\Delta))}{dT} \dot{T}. \quad (21)$$

$$\frac{1}{\tau_S} = -\frac{d \ln(VT^3)}{dT} \dot{T}. \quad (22)$$

The last term is negligible, $\tau_S \gg \tau_\Delta, \tau_T$ since pions dominate and we have near conservation of entropy which for massless particles would in fact imply $VT^3 = \text{Const.}$

Since entropy must be (slightly) increasing, while T is decreasing with time, $\tau_S > 0$. Similarly, $\tau_T > 0$, since the temperature decreases with time, and $x^2 K_2(x)$, $x = m/T$ increases with T :

$$x^2 K_2(x) \approx \sqrt{0.5\pi} x^{3/2} \exp(-x); \quad (23)$$

Therefore:

$$\frac{1}{\tau_T} \approx -\frac{m_\Delta}{T} \left(1 - \frac{3}{2} \frac{T}{m_\Delta} \dots \right) \frac{\dot{T}}{T}. \quad (24)$$

We now evaluate the magnitude of τ_T invoking a model of matter expansion of the type used e.g. in [13], where the longitudinal and transverse expansion is considered to be independent. In the proper rest frame of the out-flowing matter,

$$\frac{dS}{dy} \propto T^3 \frac{dV}{dy} = \pi R_\perp^2(\tau) T^3 \frac{dz}{dy} \simeq \text{Const.} \quad (25)$$

We will use $dz/dy \simeq \tau$, where τ is the proper time in the local volume element, this is exact for a 1-d ideal hydro flow. The growth of the transverse dimension can be generically described by:

$$R_\perp(\tau) = R_0 + \int_{\tau_0}^{\tau} v(\tau') d\tau'. \quad (26)$$

From Eq.(25) by elementary evaluation we obtain:

$$\frac{\dot{T}}{T} = -\frac{1}{3} \left(\frac{2(v\tau/R_\perp) + 1}{\tau} \right). \quad (27)$$

Equation (27) evaluated near hadronization condition is yielding the magnitude of τ_T , see Eq.(24). If the maximum expansion velocity is practically instantly achieved, $v\tau/R_\perp \simeq 1$. This leads to maximum value of $\dot{T}/T \simeq -1/\tau$. However if a more realistic profiles are assumed, \dot{T}/T is diminished in magnitude as much as 30%. We thus conclude that

$$\frac{0.5}{\tau_h} \frac{m_\Delta}{T} < \frac{1}{\tau_T} < \frac{0.7}{\tau_h} \frac{m_\Delta}{T}$$

which for hadronization time $\tau_h < 10$ fm can compete with the width of the Δ -resonance, $1/t_\Delta \simeq 120$ MeV. As this shows, the details of the expansion model are not critical for the results we obtain. In actual calculations we employ $v(\tau)$ described in [13], where we assume that the expansion is already at maximum velocity at the time of chemical freeze-out. The resulting dependence $T(\tau)$ after chemical freeze-out is shown in figure 1. We note that the time between chemical and thermal freeze-out $\Delta\tau$ is not longer than about 2.5 fm/c, and can be as short as 1 fm/c. However, even such a short scattering period is enough to alter the visible yields of strong resonances, in fact most pronounced effect we find in the latter case, since the longer time allows a greater degree of chemical equilibration.

We now can solve Eq.(20). Employing Eq.(1) we have:

$$\frac{d\Upsilon_\Delta}{d\tau} + \tilde{\Gamma}(\tau) \Upsilon_\Delta = q(\tau), \quad (28)$$

$$\tilde{\Gamma}(\tau) = \left[1 + \Upsilon_\pi \frac{N_\Delta^\infty}{N_N^\infty} \right] \frac{1}{\tau_\Delta} - \frac{1}{\tau_T}, \quad (29)$$

$$q(\tau) = \Upsilon_\pi \frac{N_0^{\text{tot}}}{N_N^\infty} \frac{1}{\tau_\Delta}, \quad (30)$$

where N_Δ^∞ and N_N^∞ are densities of Δ and N resonances with $\Upsilon_\Delta = \Upsilon_N = 1$. The solution of Eq.(28) is elementary:

$$\Upsilon_\Delta(\tau) = \left(\Upsilon_\Delta^0 + \int_{\tau_h}^{\tau} q e^{\int_{\tau_h}^{\tau'} \tilde{\Gamma} d\tau''} d\tau' \right) e^{-\int_{\tau_h}^{\tau} \tilde{\Gamma} d\tau'} \quad (31)$$

where τ_h is initial expansion time at hadronization, and $\tau_h < \tau < \tau_{\text{max}}$, upper time limit chosen to yield $T_{\text{max}} = 120$ MeV, i.e. $\tau_{\text{max}} \simeq 8$ fm.

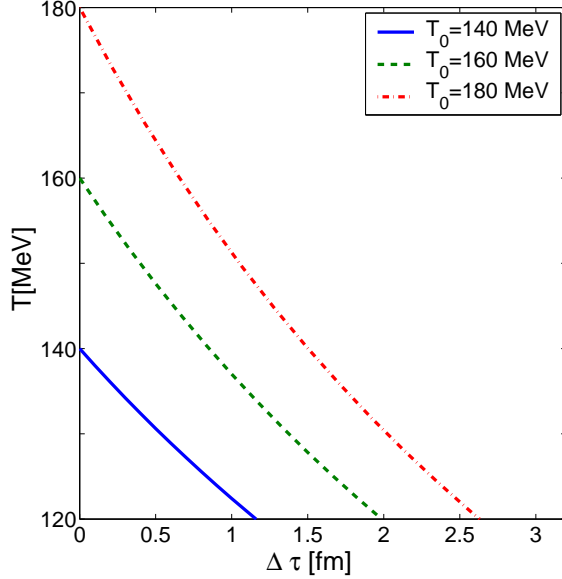


FIG. 1: Temperature T as function of $\delta\tau$, the proper time interval between chemical and thermal freeze-out or chemical freeze-out temperature (from top to bottom) $T = 180, 160, 140$ MeV and thermal freeze-out $T \geq 120$ MeV.

In order to evaluate the final Δ multiplicity we need also to know initial particles densities right after hadronization which we consider for RHIC head-on Au–Au collisions at $\sqrt{s_{NN}} = 200$ GeV. We introduce the initial hadron yields inspired by a picture of a rapid hadronization of QGP with all hadrons produced with yields governed by entropy and strangeness content of QGP by quark recombination. In this model the yields of mesons and baryons are controlled by the constituent quark fugacity γ_q :

$$\Upsilon_\pi^0 = \gamma_q^2; \quad \Upsilon_{\Delta,N}^0 = \gamma_q^3. \quad (32)$$

Thus for $\gamma_q > 1$ we have the condition $\Upsilon_\Delta < \Upsilon_\pi \Upsilon_N$, and the yield of Δ will increase in the time evolution.

For each entropy content of the QGP fireball, the corresponding fixed background value of γ_q can be found once hadronization temperature is known [12]. For $T = 140$ MeV pions form a nearly fully degenerate Bose gas with $\gamma_q \simeq 1.6$. In the following discussion, aside of this initial condition, we also consider the value pairs $T = 150$ MeV, $\gamma_q = 1.42$, $T = 160$ MeV, $\gamma_q = 1.27$, $T = 170$ MeV, $\gamma_q = 1.12$ and $T = 180$ MeV with $\gamma_q = 1$. Note further that for $m \gg T$ the density Δ is relatively low, thus there is no significant dependence of $1/\tau_\Delta$ on T and Υ_Δ ; in essence $\tau_\Delta = \hbar/\Gamma_\Delta$ takes the free space value $\tau_\Delta \simeq \hbar/120$ MeV.

In figure 2 we present results for ratios Δ/Δ_0 (solid lines) and N/N_0 (dashed lines) as functions of temperature T , beginning from the presumed initial hadronization temperature T through $T_{\max} = 120$ MeV. Δ_0 and N_0 are the initial yields obtained at each hadronization

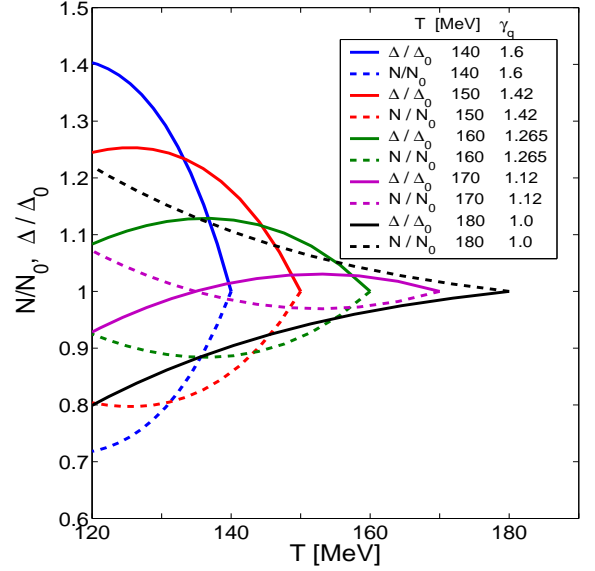


FIG. 2: The ratio Δ/Δ^0 (solid lines) and N/N^0 (dashed lines) as functions of temperature T for select given pairs of values T, γ_q , see text and figure box for details.

temperature. For $T < 180$ MeV, initially $\Upsilon_\Delta < \Upsilon_N \Upsilon_\pi$, thus based on our prior discussion, we expect that the master equation leads to an initial increase in the yield of resonances. However, as temperature drops, due to the dynamics of the expansion the increasing yield of Δ turns over, and a final nett increase of resonance yield is observed for $T \leq 160$ MeV. We note that for $T \geq 180$ MeV there is a continuous depletion of resonance yield. The nucleon yields move in opposite direction to the Δ -resonance.

This behavior can be understood in qualitative manner as follows: The total number $\Delta + N$ is conserved therefore Δ multiplicity increases and N multiplicity decreases until they reach transient chemical equilibrium ($dN_\Delta/d\tau = 0$), corresponding to the maximum point seen for Δ in figure 2. There is also influence of expansion: even if for some temperature the transient equilibrium condition (13) is reached, the system cannot stay in this equilibrium, Υ_Δ and Υ_N are increasing to conserve total number of particles. Υ_Δ increases faster because of larger Δ 's mass. After Υ_Δ becomes larger than $\Upsilon_\pi \Upsilon_N$ Δ decay begins to dominate and their multiplicity is decreasing. The special case at hadronization temperature $T = 180$ MeV where, $\Upsilon_i = 1$ and equilibrium condition is satisfied initially. As expansion sets in, Δ is decreasing because $\Upsilon_\Delta > \Upsilon_N$ (recall that here $\Upsilon_\pi = 1$). In the SHM evaluation of yields one assumes that all ratios seen in figure 2 are unity.

The initial hadronization yields which we used as reference in figure 2 are not accessible to measurement. Therefore, we consider in figure 3 the fractional yield Δ/N_{tot} (top frame), again as a function of temperature

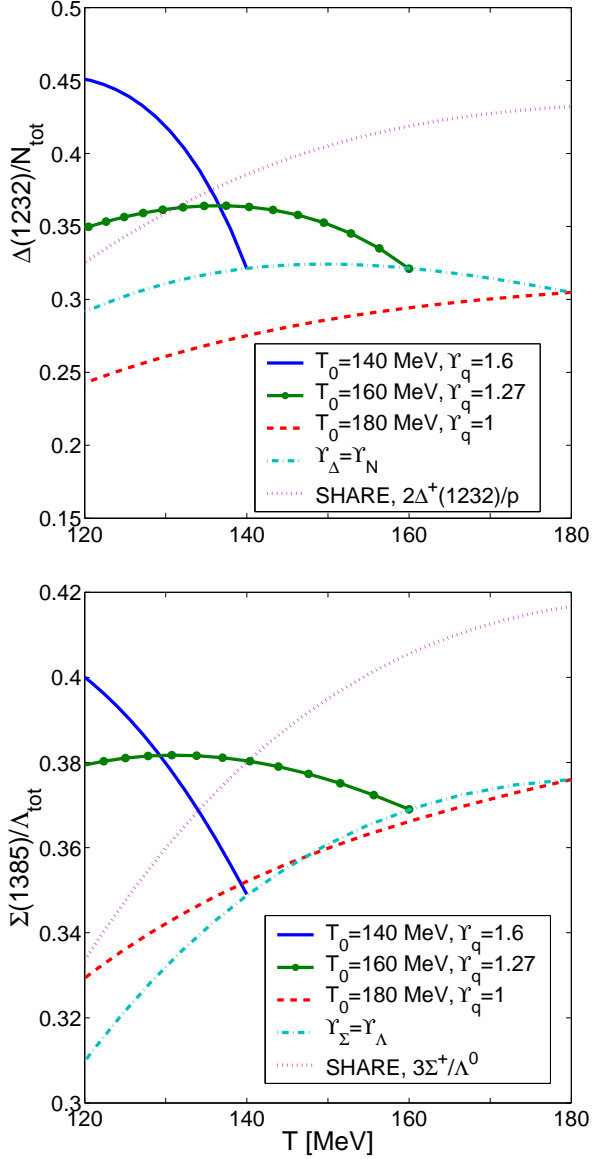


FIG. 3: Relative resonance yield, for (top) Δ/N_{tot} and (bottom) $\Sigma(1385)/\Lambda_{\text{tot}}$ as a functions of freeze-out temperature, for hadronization temperatures $T_0 = 140, 160, 180$ MeV, see box and text for details. The dotted brown line gives the expected SHM chemical equilibrium result.

T . The results for hadronization temperatures $T_0 = 140$ (solid blue line), $T_0 = 160$ (dash-dot green line) and $T_0 = 180$ MeV (dashed brown line) are shown. N_{tot} is fixed by hadronization condition and is not a function of time, as discussed. Thus the observable final rapidity nucleon yield corresponds to the initial value at hadronization. Note that up to strange and multi strange baryon contribution, N_{tot} is the total baryon (rapidity) yield.

Since in this study we have considered a subset of all relevant baryon resonances our chemical equilibrium reference yield (line for $\Upsilon_\Delta = \Upsilon_N$) is not the same

as the corresponding reference line for the full statistical hadronization model (SHM) evaluation, obtained using SHARE2, and presented as $2\Delta^+/p$ (upper frame) and $3\Sigma^+/\Lambda^0$ (lower frame). The SHARE2-SHM value $\Delta^{++}/p \simeq 0.2$ at $T \simeq 160$ MeV is consistent with the STAR d-Au results [14]. Also, comparing our with the SHARE2 result we note that SHARE2 yield is larger at chemical freeze-out. The magnitude of the difference in the yields at time of chemical freeze-out provides a measure of the magnitude of the corrections we can expect to arise in the full treatment at thermal freeze-out and/or systematic error for these yields.

The nature of these effects is different for the two yield cases considered: the presence of heavier resonances which cascade by way of Δ leads to an increase of the thermal freeze-out yield. The correction is thus nearly as much as we see the SHARE2 yield higher at chemical freeze-out. For $\Sigma(1385)$ the difference with SHARE2 arises from a difference of contributions of partial decays producing Λ_{tot} , thus the correction is multiplicative factor which does not change, but is uncertain in magnitude due to lack of knowledge about the branching ratios.

We believe that the Δ and $\Sigma(1385)$ yields are underestimated by about 15% – 35%. (bigger effect for hadronization at higher T). This implies that depending on hadronization temperature a relative yield range $0.16 < \Delta^{++}/p = 0.5\Delta/N_{\text{tot}} < 0.26$ arises, and similarly (see lower frame in figure 3) $0.35 < \Sigma(1385)/\Lambda_{\text{tot}} < 0.43$ with the *higher* relative yield corresponding to the *lower* hadronization temperature. One of the key results of this work is the narrow range for $\Sigma(1385)/\Lambda_{\text{tot}}$, and the fact that the initial chemical non-equilibrium effect leads to a reversal of the SHM model situation: the relative yields of massive resonances decreases with decreasing hadronization temperature. .

In order to compare with the experimental results we note that the data presented [1, 2] are for charged $\Sigma(1385)$, particle and antiparticle channels, $(\Sigma^\pm(1385) + \overline{\Sigma^\pm(1385)})/(\Lambda_{\text{tot}} + \overline{\Lambda_{\text{tot}}}) \simeq 0.29$. This result needs to be multiplied with $3/2$ to be comparable to results presented here which include $\Sigma^0(1385)$. Multiplying the value for hadronization at $T = 140$ MeV with thermal freeze-out at $T = 120$ MeV, and allowing for contribution by heavier resonances as indicated by SHARE2 our result is in perfect agreement with [1, 2] However, given the narrow range of results we find, it seems that the high yield of $\Sigma(1385)$, seen the error $\mathcal{O}(20\%)$ is nearly compatible with the entire range of chemical freeze-out temperatures here considered – the low T chemical freeze-out is favored by 1.5 s.d. over high T .

The reader should take note that the ‘thermal’ model result presented in Ref. [1] corresponds to initial high temperature freeze-out in chemical equilibrium which is unobservable, since the high T hadronization resonance decay products have no chance to escape into free space. Thus this comparison of this model with experiment is flawed. The evolved yield is shown as (red) dashed line in figure 3, and is found 25% below the value measured. The

reason this happens is that the high T chemical freeze-out happens near chemical equilibrium and the yields follow closely the chemical equilibrium yield described by temperature, thus it is the *thermal freeze out temperature which in this case controls the final observable resonance yield*.

In summary, we have presented master equation governing the evolution in time of the $\Delta, \Sigma(1385)$ baryon resonance yield after QGP hadronization, allowing for resonance decay and production process. We have shown considering the properties of the master equation that if the yield of hadrons is initially above chemical equilibrium, the resonance population increases beyond the initial yield. Conversely, we find that in a physical system in which the particle multiplicities of hadrons arise below chemical equilibrium yields, a circumstance expected below threshold to QGP formation, the final yield of resonances is suppressed by the dominance of the resonance decay process over back reaction resonance production.

In a quantitative model we evolved the yields after QGP hadronization allowing for initial chemical non-equilibrium particle abundances, and volume expansion assuring entropy conservation. We found, see figure 3, that the thermal freeze-out fractional resonance yield differs significantly from the chemical-freeze out SHM expectation, with the scenario involving high- T hadronization resonance yield being depleted, and low- T hadronization yield scenario further enhanced in relation to the total yield.

The resonance enhancement effect we presented can only occur when the initial state is out of chemical equilibrium, and the decay/formation processes are fast enough to compete with the hadron volume evolution. One would thus think that ‘narrow’, i.e. quasi-stable resonances are not subject to the effects considered here. However, a special consideration must be given to nar-

row resonance which are strongly coupled to more massive resonances which can decay fast into other channels. An example is $\Lambda(1520)$. In this situation the lower mass state, here $\Lambda(1520)$, is the ground state which is depleted by coupling to the yet more massive state. There are several states of relevance to consider, hence $\Lambda(1520)$ ‘quenching’ [7] must be addressed in a more rigorous numerical approach. We will to return to discuss this case in a separate publication [15].

Aside of several specific prediction we made here, there are three important general consequences of our study: a) the fractional yield of resonances A^*/A can be considerably higher than expected naively in SHM model of QGP hadronization, b) since there is nearly a factor two difference in the final thermal freeze-out ratio in Δ/N_{tot} , while the SHM yields a more T independent result, one can imagine the use of Δ/N_{tot} as a tool to distinguish the different hadronization conditions e.g. chemical non-equilibrium vs chemical equilibrium a point noted in similar context before [16]; and c) we have shown that the relatively high yield of charged $\Sigma^\pm(1385)$ reported by STAR is well explained by our considerations with hadronization at $T = 140$ MeV being favored.

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